# Flavor violation via Planck scale alignment in the 2HDM


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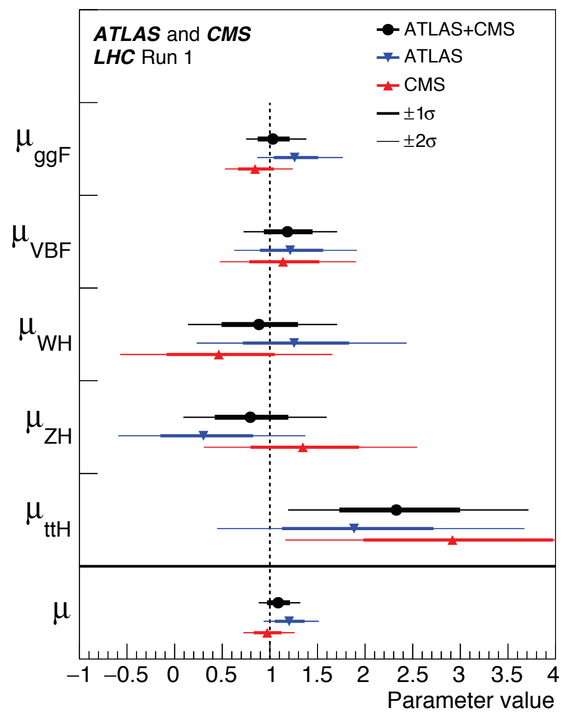
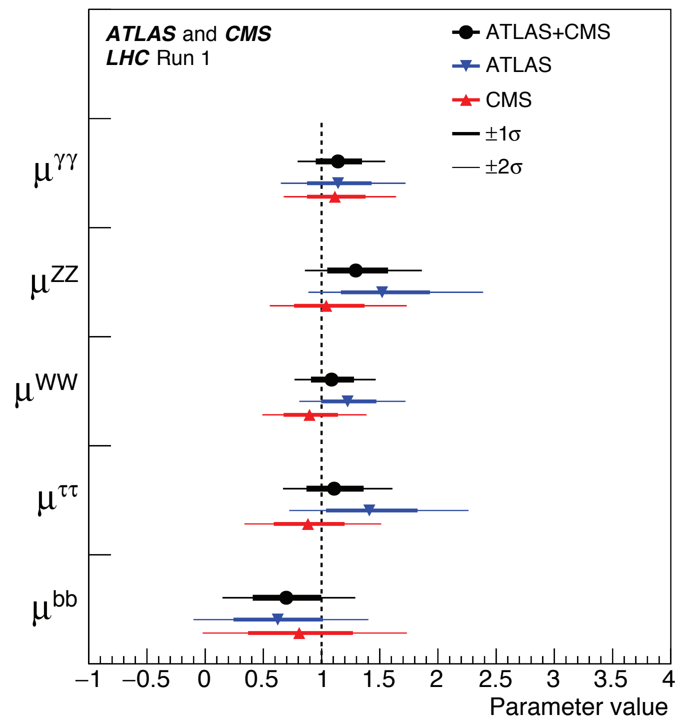
This talk is based on work in collaboration with Stefania Gori and Edward Santos. Coming soon to an arXiv near you.

# Introduction

The Standard Model (SM) remains a surprisingly accurate description of particle physics at the TeV scale. The properties of the observed Higgs boson remain consistent with SM predictions (given the statistical power of the Higgs data).




Reference: G. Aad *et al.* [ATLAS and CMS Collaborations], JHEP **1608**, 045 (2016) [arXiv:1606.02266 [hep-ex]].

So, why are we having this conference, entitled ”Workshop on Multi-Higgs Models?”*

In fact, by the end of this conference, you will have plenty of motivations for why we are interested in non-minimal Higgs sectors. But, apart from all such motivations, consider the following.

Given that fermionic matter of the SM is non-minimal why shouldn't scalar matter also be non-minimal? (To paraphrase I.I. Rabi, “who ordered that?”). In my opinion, one of the most important questions that the LHC can answer is: are there additional Higgs bosons to be discovered (at the TeV scale)?

*What's worse is that there is not even a cool acronym to impress our friends!

Electroweak data already imposes strong constraints on possible Higgs sector extensions.

1. The electroweak $\rho$-parameter, $\rho \equiv m_W^2/(m_Z^2 \cos^2\theta_W) \simeq 1$ strongly suggests that extended Higgs sectors should contain at most only scalar doublets and singlets.[†]
2. Generic Yukawa couplings of an extended Higgs sector yield tree-level Higgs-mediated flavor changing neutral currents (FCNCs) at a level far greater than that which can be tolerated in light of flavor physics data.

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[†]For general scalar multiplets, one typically achieves $\rho \simeq 1$ by an unnatural fine-tuning of the Higgs scalar potential. Even in the Georgi-Macacek model which contains both scalar doublet and triplets with a custodial symmetric scalar potential, one finds that the custodial symmetric form of the potential is not stable under radiative corrections.

# FCNCs and the two-Higgs doublet model (2HDM)

Henceforth, we consider the two-Higgs-doublet extension of the SM. The 2HDM Higgs-quark Yukawa Lagrangian (in terms of quark mass-eigenstates) is:

$$-\mathscr{L}_{\rm Y} = \overline{U}_L \Phi_i^{0\,*} h_i^U U_R - \overline{D}_L K^\dagger \Phi_i^- h_i^U U_R + \overline{U}_L K \Phi_i^+ h_i^{D\,\dagger} D_R + \overline{D}_L \Phi_i^0 h_i^{D\,\dagger} D_R + {\rm h.c.}\,,$$

where $K$ is the CKM mixing matrix, and there is an implicit sum over the two Higgs fields ($i = 1, 2$). The $h^{U,D}$ are $3 \times 3$ Yukawa coupling matrices.

In order to naturally eliminate tree-level Higgs-mediated FCNC, we shall impose a discrete symmetry $\Phi_1 \to +\Phi_1$ and $\Phi_2 \to -\Phi_2$ to restrict the structure of $\mathscr{L}_{\rm Y}$. Two different choices for how the discrete symmetry acts on the quarks then yield:

- Type-I Yukawa couplings: $h_1^U = h_1^D = 0$,
- Type-II Yukawa couplings: $h_1^U = h_2^D = 0$.

For simplicity in the presentation below, assume that the Higgs scalar potential and vacuum are CP-invariant. In the $\Phi_1$–$\Phi_2$ basis, we define $\tan\beta \equiv v_2/v_1$ and $\alpha$ as the angle that diagonalizes the CP-even Higgs squared-mass matrix. Then, the neutral Higgs interactions are

$$-\mathscr{L}_Y = \frac{1}{v}\sum_{F=U,D,E}\overline{F}\left\{s_{\beta-\alpha}M_F + c_{\beta-\alpha}M_F^{1/2}\left[\rho_R^F + i\varepsilon_F\gamma_5\rho_I^F\right]M_F^{1/2}\right\}Fh$$

$$+\frac{1}{v}\sum_{F=U,D,E}\overline{F}\left\{c_{\beta-\alpha}M_F - s_{\beta-\alpha}M_F^{1/2}\left[\rho_R^F + i\varepsilon_F\gamma_5\rho_I^F\right]M_F^{1/2}\right\}FH$$

$$+\frac{1}{v}\sum_{F=U,D,E}\overline{F}\left\{M_F^{1/2}\left(\rho_I^F - i\varepsilon_F\gamma_5\rho_R^F\right)M_F^{1/2}\right\}FA$$

where $s_{\beta-\alpha} \equiv \sin(\beta-\alpha)$, $c_{\beta-\alpha} \equiv \cos(\beta-\alpha)$, and

$$\varepsilon_F = \begin{cases} +1 & \text{for } F = U\,, \\ -1 & \text{for } F = D, E\,. \end{cases}$$

Note that $M_F$ are the diagonal fermion matrices (neutrinos are assumed massless) and the $\rho^F_{R,I}$ are arbitrary $3\times 3$ Hermitian matrices that are in general non-diagonal in generation space. Hence, tree-level FCNCs mediated by neutral Higgs bosons are present (as well as new sources of CP-violation).

## Definitions of $\boldsymbol{\rho^F_{R,I}}$

$$M_F^{1/2}\rho^F_R M_F^{1/2}=\frac{v}{2\sqrt{2}}(\rho^F+[\rho^F]^\dagger)\,,\qquad iM_F^{1/2}\rho^F_I M_F^{1/2}=\frac{v}{2\sqrt{2}}(\rho^F-[\rho^F]^\dagger)\,,$$

where $\rho^F\equiv\epsilon_{ij}h^F_j v_i/v$ with $\langle\Phi^0_i\rangle\equiv v_i/\sqrt{2}$ and $v^2\equiv v_1^2+v_2^2=(246~\text{GeV})^2$. We can define an analogous quantity, $\kappa^F\equiv\sqrt{2}M_F/v=h^F_i v^*_i/v$. Note that $\kappa^F$ is proportional to the diagonal fermion mass matrix.

Remark: $\kappa^F$ and $\rho^F$ are Higgs-fermion Yukawa matrices in the Higgs basis.

In the CP-conserving Type-I and Type-II 2HDM, $\rho^{I,D}_I=0$ and[‡]

$$\text{Type I}:\qquad \rho^D_R=\rho^U_R=\mathbb{1}\cot\beta\,,$$

$$\text{Type II}:\qquad \rho^D_R=-\mathbb{1}\tan\beta\,,\qquad \rho^U_R=\mathbb{1}\cot\beta\,,$$

where $\mathbb{1}$ is the $3\times 3$ identity matrix. Thus, the neutral Higgs-fermion couplings are flavor diagonal!

[‡]In Type-I and Type-II models, the couplings to leptons follows the pattern of the down-type quark couplings. In the so-called Types Y and X models, the Types I and II quark couplings are associated with Types II and I lepton couplings, respectively.

# The flavor-aligned two-Higgs doublet model (A2HDM)

We can by fiat declare that $\rho^F = a^F \kappa^F$ for $F = U, D, E$, were $a^F$ is called the alignment parameter.[§] It follows that

$$\rho_R^F = (\text{Re } a^F)\mathbb{1}\,, \qquad\qquad \rho_I^F = (\text{Im } a^F)\mathbb{1}\,.$$

The corresponding neutral Higgs–fermion Yukawa couplings are given by

$$\begin{aligned}
-\mathscr{L}_Y = &\frac{1}{v}\sum_{F=U,D,E} \overline{F} M_F \Big\{ s_{\beta-\alpha} + c_{\beta-\alpha}\big[\text{Re } a^F + i\epsilon^F \text{Im } a^F \gamma_5\big]\Big\} F h \\
&+ \frac{1}{v}\sum_{F=U,D,E} \overline{F} M_F \Big\{ c_{\beta-\alpha} - s_{\beta-\alpha}\big[\text{Re } a^F + i\epsilon^F \text{Im } a^F \gamma_5\big]\Big\} F H \\
&+ \frac{1}{v}\sum_{F=U,D,E} \overline{F} M_F \Big\{ \big[\text{Im } a^F - i\epsilon^F \text{Re } a^F \gamma_5\big]\Big\} F A\,,
\end{aligned}$$

and the Higgs-fermion couplings are diagonal as advertised.[¶]

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[§] A. Pich and P. Tuzon, Phys. Rev. D **80**, 091702 (2009) [arXiv:0908.1554 [hep-ph]].

[¶] In the Types I, II X and Y 2HDMs, the alignment parameters are fixed to either $\cot\beta$ or $-\tan\beta$.

## Radiative stability of the flavor aligned 2HDM

The flavor-alignment conditions of the A2HDM are not radiatively stable, except in the case of the Types I, II X and Y 2HDMs. Indeed, as shown by P.M. Ferreira, L. Lavoura and J.P. Silva, Phys. Lett. B **688**, 341 (2010) [arXiv:1001.2561 [hep-ph]], flavor alignment is preserved by the renormalization-group (RG) running of the Yukawa coupling matrices only in the cases of the standard type-I, II, X, and Y models. This means that the A2HDM is an artificially tuned model.

Our proposal is to examine the possibility that the flavor alignment condition is imposed at the Planck scale,[‖] due to new physics that is presently unknown. One can then use an RG analysis to determine the structure of the Higgs-fermion Yukawa couplings at the electroweak scale. This in turn will lead to small flavor-violation in the neutral Higgs-quark interactions that can be constrained by current and future experiments.

[‖]This ansatz was first considered by C.B. Braeuninger, A. Ibarra and C. Simonetto, Phys. Lett. B **692**, 189 (2010) [arXiv:1005.5706 [hep-ph]].

# RG equations for the Yukawa coupling matrices

Prior to diagonalizing the fermion mass matrices, we define Yukawa coupling matrices $\eta_a^{F,0}$, for $a = 1, 2$ and $F = U$, $D$ and $E$. Defining $\mathcal{D} \equiv 16\pi^2\mu(d/d\mu)$, the RGEs are given by (Ferreira, Lavoura and Silva, op. cit.),

$$\mathcal{D}\eta_a^{U,0} = -\left(8g_s^2 + \tfrac{9}{4}g^2 + \tfrac{17}{12}g'^2\right)\eta_a^{U,0} + \left\{3\mathrm{Tr}\left[\eta_a^{U,0}(\eta_{\bar{b}}^{U,0})^\dagger + \eta_a^{D,0}(\eta_{\bar{b}}^{D,0})^\dagger\right] + \mathrm{Tr}\left[\eta_a^{E,0}(\eta_{\bar{b}}^{E,0})^\dagger\right]\right\}\eta_b^{U,0}$$

$$-2(\eta_{\bar{b}}^{D,0})^\dagger\eta_a^{D,0}\eta_b^{U,0} + \eta_a^{U,0}(\eta_{\bar{b}}^{U,0})^\dagger\eta_b^{U,0} + \tfrac{1}{2}(\eta_{\bar{b}}^{D,0})^\dagger\eta_b^{D,0}\eta_a^{U,0} + \tfrac{1}{2}\eta_b^{U,0}(\eta_{\bar{b}}^{U,0})^\dagger\eta_a^{U,0}\,,$$

$$\mathcal{D}\eta_a^{D,0} = -\left(8g_s^2 + \tfrac{9}{4}g^2 + \tfrac{5}{12}g'^2\right)\eta_a^{D,0} + \left\{3\mathrm{Tr}\left[(\eta_{\bar{b}}^{D,0})^\dagger\eta_a^{D,0} + (\eta_{\bar{b}}^{U,0})^\dagger\eta_a^{U,0}\right] + \mathrm{Tr}\left[(\eta_{\bar{b}}^{E,0})^\dagger\eta_a^{E,0}\right]\right\}\eta_b^{D,0}$$

$$-2\eta_b^{D,0}\eta_a^{U,0}(\eta_{\bar{b}}^{U,0})^\dagger + \eta_b^{D,0}(\eta_{\bar{b}}^{D,0})^\dagger\eta_a^{D,0} + \tfrac{1}{2}\eta_a^{D,0}\eta_b^{U,0}(\eta_{\bar{b}}^{U,0})^\dagger + \tfrac{1}{2}\eta_a^{D,0}(\eta_{\bar{b}}^{D,0})^\dagger\eta_b^{D,0}\,,$$

$$\mathcal{D}\eta_a^{E,0} = -\left(\tfrac{9}{4}g^2 + \tfrac{15}{4}g'^2\right)\eta_a^{E,0} + \left\{3\mathrm{Tr}\left[(\eta_{\bar{b}}^{D,0})^\dagger\eta_a^{D,0} + (\eta_{\bar{b}}^{U,0})^\dagger\eta_a^{U,0}\right] + \mathrm{Tr}\left[(\eta_{\bar{b}}^{E,0})^\dagger\eta_a^{E,0}\right]\right\}\eta_b^{E,0}$$

$$+\eta_b^{E,0}(\eta_{\bar{b}}^{E,0})^\dagger\eta_a^{E,0} + \tfrac{1}{2}\eta_a^{E,0}(\eta_{\bar{b}}^{E,0})^\dagger\eta_b^{E,0}\,.$$

These equations take the form in any basis of scalar fields. Applying these results to the Higgs basis yields the RGEs for the $\kappa^{F,0}$ and $\rho^{F,0}$.

We now identify the fermion mass eigenstates,

$$P_L U = V_L^U P_L U^0\,, \quad P_R U = V_R^U P_R U^0\,, \quad P_L D = V_L^D P_L D^0\,, \quad P_R D = V_R^D P_R D^0\,,$$
$$P_L E = V_L^E P_L E^0\,, \quad P_R E = V_R^D P_R E^0\,, \quad P_L N = V_L^E P_L N^0\,,$$

and the Cabibbo-Kobayashi-Maskawa (CKM) matrix is defined as $K \equiv V_L^U V_L^{D\,\dagger}$. Note that for the neutrino fields, we are free to choose $V_L^N = V_L^E$ since neutrinos are exactly massless in this analysis.

In particular, the unitary matrices $V_L^F$ and $V_R^F$ (for $F = U$, $D$ and $E$) are chosen such that

$$M_U = \frac{v}{\sqrt{2}} V_L^U \kappa^{U,0} V_R^{U\,\dagger} = \mathrm{diag}(m_u\,,\, m_c\,,\, m_t)\,,$$

$$M_D = \frac{v}{\sqrt{2}} V_L^D \kappa^{D,0\,\dagger} V_R^{D\,\dagger} = \mathrm{diag}(m_d\,,\, m_s\,,\, m_b)\,,$$

$$M_E = \frac{v}{\sqrt{2}} V_L^E \kappa^{E,0\,\dagger} V_R^{E\,\dagger} = \mathrm{diag}(m_e\,,\, m_\mu\,,\, m_\tau)\,.$$

The $\kappa^F$ and $\rho^F$ matrices previously defined are given by,

$$\kappa^U = V_L^U \kappa^{U,0} V_R^{U\,\dagger}\,, \qquad \kappa^D = V_R^D \kappa^{D,0} V_L^{D\,\dagger}\,, \qquad \kappa^E = V_R^D \kappa^{E,0} V_L^{E\,\dagger}\,,$$
$$\rho^U = V_L^U \rho^{U,0} V_R^{U\,\dagger}\,, \qquad \rho^D = V_R^D \rho^{D,0} V_L^{D\,\dagger}\,, \qquad \rho^E = V_R^D \rho^{E,0} V_L^{E\,\dagger}\,.$$

We can therefore obtain the RGEs for $\kappa^F$ and $\rho^F$. In this analysis, the diagonalization of the fermion mass matrices are carried out at the electroweak scale.[**] As a result, the $V_L^F$ and $V_R^F$ are fixed matrices and the CKM matrix $K$ does not run. Only Yukawa couplings evolve under RG running. That is, the running Yukawa couplings are defined with respect to a fixed fermion basis.

The end result is that the RGEs for the $\kappa^F$ and $\rho^F$ explicitly contain factors of the CKM matrix $K$. Thus, if $\kappa^F$ and $\rho^F$ are proportional at one energy scale, they will no longer be proportional at another scale.

[**] In practice, one should consider carefully how Yukawa couplings evolve from the electroweak scale down to the scale at which the corresponding pole masses are defined. We neglect these effects, as they are numerically small.

For example,

$$
\begin{aligned}
\mathcal{D}\kappa^U &= -\left(8g_s^2 + \tfrac{9}{4}g^2 + \tfrac{17}{12}g'^{\,2}\right)\kappa^U + \left\{3\mathrm{Tr}\left[\kappa^U\kappa^{U\,\dagger} + \kappa^D\kappa^{D\,\dagger}\right] + \mathrm{Tr}\left[\kappa^E\kappa^{E\,\dagger}\right]\right\}\kappa^U \\
&\quad + \left\{3\mathrm{Tr}\left[\kappa^U\rho^{U\,\dagger} + \kappa^D\rho^{D\,\dagger}\right] + \mathrm{Tr}\left[\kappa^E\rho^{E\,\dagger}\right]\right\}\rho^U - 2K\left(\kappa^{D\,\dagger}\kappa^D K^\dagger\kappa^U + \rho^{D\,\dagger}\kappa^D K^\dagger\rho^U\right) \\
&\quad + \kappa^U(\kappa^{U\,\dagger}\kappa^U + \rho^{U\,\dagger}\rho^U) + \tfrac{1}{2}K(\kappa^{D\,\dagger}\kappa^D + \rho^{D\,\dagger}\rho^D)K^\dagger\kappa^U \\
&\quad + \tfrac{1}{2}(\kappa^U\kappa^{U\,\dagger} + \rho^U\rho^{U\,\dagger})\kappa^U\,,
\end{aligned}
$$

$$
\begin{aligned}
\mathcal{D}\rho^U &= -\left(8g_s^2 + \tfrac{9}{4}g^2 + \tfrac{17}{12}g'^{\,2}\right)\rho^U + \left\{3\mathrm{Tr}\left[\rho^U\kappa^{U\,\dagger} + \rho^D\kappa^{D\,\dagger}\right] + \mathrm{Tr}\left[\rho^E\kappa^{E\,\dagger}\right]\right\}\kappa^U \\
&\quad + \left\{3\mathrm{Tr}\left[\rho^U\rho^{U\,\dagger} + \rho^D\rho^{D\,\dagger}\right] + \mathrm{Tr}\left[\rho^E\rho^{E\,\dagger}\right]\right\}\rho^U - 2K\left(\kappa^{D\,\dagger}\rho^D K^\dagger\kappa^U + \rho^{D\,\dagger}\rho^D K^\dagger\rho^U\right) \\
&\quad + \rho^U(\kappa^{U\,\dagger}\kappa^U + \rho^{U\,\dagger}\rho^U) + \tfrac{1}{2}K(\kappa^{D\,\dagger}\kappa^D + \rho^{D\,\dagger}\rho^D)K^\dagger\rho^U \\
&\quad + \tfrac{1}{2}(\kappa^U\kappa^{U\,\dagger} + \rho^U\rho^{U\,\dagger})\rho^U\,,
\end{aligned}
$$

# Flavor-aligned Yukawa coupling matrices at $\Lambda = M_{\rm PL}$

Our setup is as follows. We assume flavor-alignment at the Planck scale, $\Lambda = M_{\rm PL}$,

$$\rho^Q(\Lambda) = a^Q \kappa^Q(\Lambda), .$$

We assume that there exists a low-energy scale $\Lambda_H$ that characterizes the mass scale of the second Higgs doublet. We take $\Lambda_H > 400$ GeV, in order that the observed Higgs boson possess SM-like properties (within about 20%). This corresponds to the decoupling limit. To be consistent with the observe quark masses and CKM matrix, we impose

$$\kappa^Q(\Lambda_H) = \sqrt{2} M_Q(\Lambda_H)/v \, .$$

where the $M_Q$ $(Q = U,\, D)$ are the diagonal quark matrices. We therefore have two boundary conditions, one at the high scale and one at the low scale.

We begin by assuming flavor-alignment at $\Lambda_H$ via a low-scale alignment parameter $a'^Q$ in the first approximation of an iterative process, $\rho^Q(\Lambda_H) = a'^Q \kappa^Q(\Lambda_H)$. We then decompose $\rho^Q(\Lambda)$ into parts that are aligned and misaligned with $\kappa^Q(\Lambda)$, respectively,

$$\rho^Q(\Lambda) = a^Q \kappa^Q(\Lambda) + \delta\rho^Q,$$

where $a^Q$ represents the aligned part (in general, different from $a'^Q$), and $\delta\rho^Q$ the corresponding degree of misalignment at the high scale.

To minimize the misaligned part of $\rho^Q(\Lambda)$, we implement the cost function,

$$\Delta^Q \equiv \sum_{i,j=1}^{3} |\delta\rho^Q_{ij}|^2 = \sum_{i,j=1}^{3} |\rho^Q_{ij}(\Lambda) - a^Q \kappa^Q_{ij}(\Lambda)|^2,$$

which once minimized, provides the optimal value of the complex parameter $a^Q$ for flavor-alignment at the high scale,

$$a^Q \equiv \frac{\sum_{i,j=1}^{3} \kappa^{Q*}_{ij}(\Lambda)\rho^Q_{ij}(\Lambda)}{\sum_{i,j=1}^{3} \kappa^{Q*}_{ij}(\Lambda)\kappa^Q_{ij}(\Lambda)}.$$

We subsequently impose flavor-alignment at the high scale using this optimized alignment parameter,

$$\rho^Q(\Lambda) = a^Q \kappa^Q(\Lambda),$$

and evolve the one-loop RGEs back down to $\Lambda_H$. At $\Lambda_H$, we match the boundary conditions for the 2HDM and SM. At this point, the matrices $\kappa^U$ and $\kappa^D$ at the scale $\Lambda_H$ are no longer diagonal, so we must rediagonalize $\kappa^U$ and $\kappa^D$ [while respectively transforming $\rho^U$ and $\rho^D$ (at the scale $\Lambda_H$)].

One can now evolve $\kappa^U$ and $\kappa^D$ down to the electroweak scale using the one-loop SM RGEs. If any of the quark masses differ from their experimental values by more than 3%, we reestablish the correct quark masses at the electroweak scale, run back up to $\Lambda_H$, and then rerun this procedure repeatedly until the two boundary conditions are satisfied. The result is flavor-alignment between $\kappa^Q(\Lambda)$ and $\rho^Q(\Lambda)$, and a set of $\rho^Q$ matrices at the electroweak scale that provide a source of FCNCs.

## The one-loop leading logarithmic approximation

$$\rho^U(\Lambda_H) \sim a^U \kappa^U(\Lambda_H) + \frac{1}{16\pi^2} \log\left(\frac{\Lambda_H}{\Lambda}\right)(\mathcal{D}\rho^U - a^U \mathcal{D}\kappa^U),$$

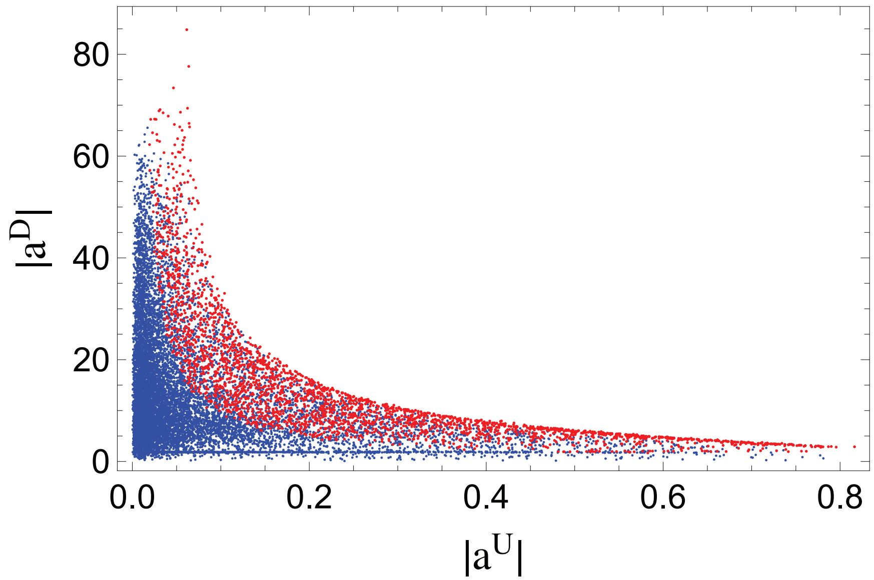
$$\rho^D(\Lambda_H) \sim a^D \kappa^D(\Lambda_H) + \frac{1}{16\pi^2} \log\left(\frac{\Lambda_H}{\Lambda}\right)(\mathcal{D}\rho^D - a^D \mathcal{D}\kappa^D)\,.$$

where $\kappa^U(\Lambda_H)$ and $\kappa^D(\Lambda_H)$ are proportional to the diagonal quark mass matrices, $M_U$ and $M_D$ respectively, at the scale $\Lambda_H$. Working to one loop order and neglecting higher order terms,

$$\rho^U(\Lambda_H)_{ij} \simeq a^U \delta_{ij} \frac{\sqrt{2}(M_U)_{jj}}{v} + \frac{(M_U)_{jj}}{4\sqrt{2}\pi^2 v^3} \log\left(\frac{\Lambda_H}{\Lambda}\right) \Big\{(a^E - a^U)[1 + a^U (a^E)^*]\delta_{ij} \sum_k (M_E^2)_{kk}$$
$$+(a^D - a^U)[1 + a^U (a^D)^*] \sum_k [3\delta_{ij}(M_D^2)_{kk} - 2(M_D^2)_{kk} K_{ik} K^*_{jk}]\Big\}\,,$$

$$\rho^D(\Lambda_H)_{ij} \simeq a^D \delta_{ij} \frac{\sqrt{2}(M_D)_{ii}}{v} + \frac{(M_D)_{ii}}{4\sqrt{2}\pi^2 v^3} \log\left(\frac{\Lambda_H}{\Lambda}\right) \Big\{(a^E - a^D)[1 + a^D (a^E)^*]\delta_{ij} \sum_k (M_E^2)_{kk}$$
$$+(a^U - a^D)[1 + a^D (a^U)^*] \sum_k [3\delta_{ij}(M_U^2)_{kk} - 2(M_U^2)_{kk} K^*_{ki} K_{kj}]\Big\}\,.$$

The validity of the one-loop leading log approximation breaks down for large values of the alignment parameters.




Blue: region of the A2HDM parameter space where the prediction for all the off-diagonal terms of the $\rho^Q$ matrices lies within a factor of 3 from the results obtained with the full running. Red: region where the one-loop leading log approximation differs significantly from the the results obtained by numerically solving the RGEs.

Remark: In our numerical analysis, we require that no Landau pole singularities appear below $\Lambda = M_{\rm PL}$. This constraint is reflected in the upper boundary of the red curve shown above.

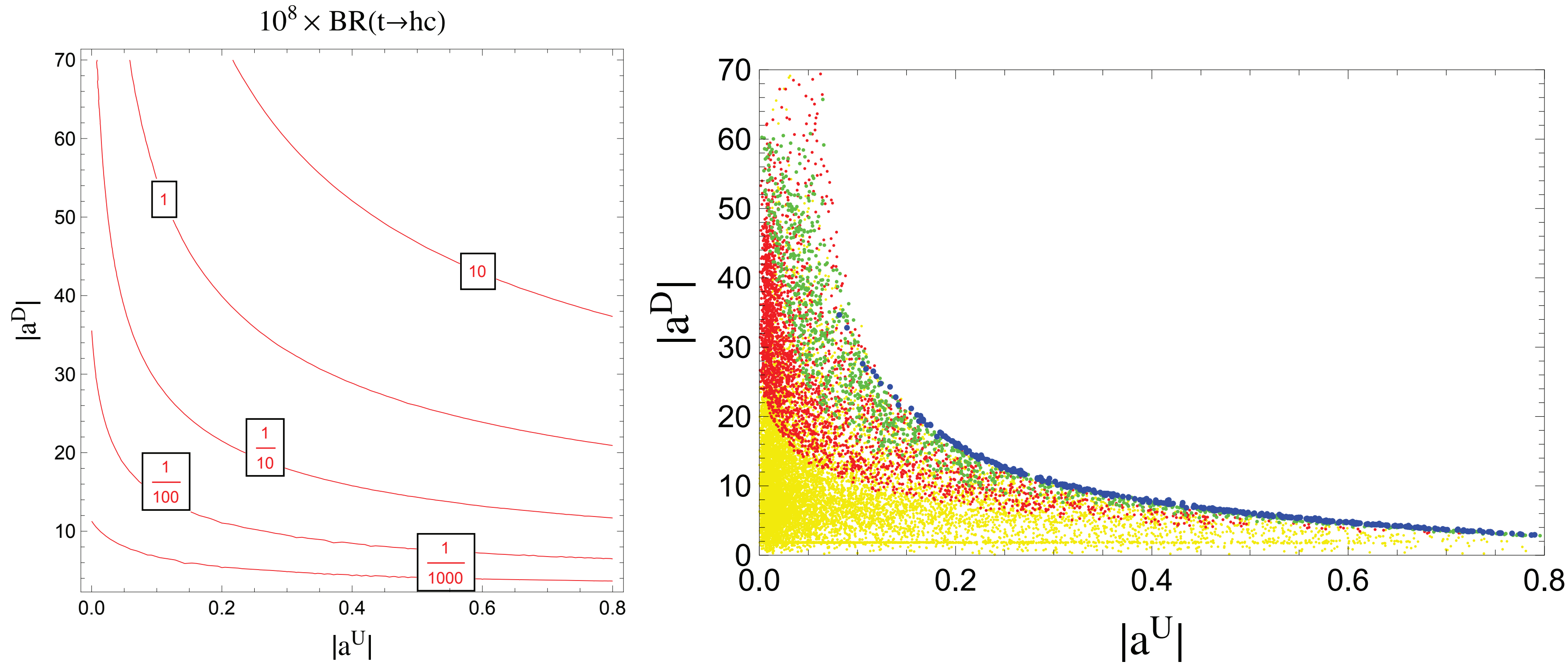
# Phenomenological consequences

1. Flavor-changing top decays.

$$\mathrm{BR}(t \to u_i h) = \cos^2(\beta-\alpha)(|\rho^U_{i3}|^2+|\rho^U_{3i}|^2)\times\frac{v^2}{4m_t^2}\frac{(1-m_h^2/m_t^2)^2}{(1-m_W^2/m_t^2)^2(1+2m_W^2/m_t^2)}\eta_{QCD}\,,$$

where $\eta_{QCD} = 1 + 0.97\alpha_s \sim 1.10$ is the NLO QCD correction to the branching ratio.

Remark: In the SM, $\mathrm{BR}(t \to ch) \sim 3 \times 10^{-15}$. Projections for the HL-LHC show that the bounds on the branching ratios of flavor violating top decays will likely be at the $10^{-4}$ level. At a future 100 TeV proton-proton machine with a large luminosity, recent estimates suggest that branching ratios as small as $\sim 10^{-7}$ could be probed with 10 ab$^{-1}$ luminosity.




Left: we use the leading log approximation to obtain $10^8\times$ BR$(t \to ch)$ . Right: the same but scanning the parameter space. Yellow, red, green and blue colors correspond to branching ratios $< 10^{-11}$, $[10^{-11} - 10^{-10}]$, $[10^{-10} - 10^{-8}]$, $> 10^{-8}$. We have fixed $\beta - \alpha = \pi/2 - 0.2$ and $\Lambda_H = 400$ GeV.

2. $B_{s,d} \to \mu^+\mu^-$.

For our calculations, we use[††]

$$\frac{\mathrm{BR}(B_{s,d} \to \mu^+\mu^-)}{\mathrm{BR}(B_{s,d} \to \mu^+\mu^-)_{\mathrm{SM}}} \simeq \left(|S_{s,d}|^2 + |P_{s,d}|^2\right) \times \left(1 + y_{s,d} \frac{\mathrm{Re}(P_{s,d}^2) - \mathrm{Re}(S_{s,d}^2)}{|S_{s,d}|^2 + |P_{s,d}|^2}\right) \left(\frac{1}{1 + y_{s,d}}\right).$$

Above, $\mathrm{BR}(B_{s,d} \to \mu^+\mu^-)_{\mathrm{SM}}$ is the SM prediction for the branching ratio extracted from an untagged rate, $y_s = (8.8 \pm 1.4)\%$ and $y_d = 0$, and

$$S_{s,d} \equiv \frac{m_{B_{s,d}}}{2m_\mu} \frac{(C_{s,d}^S - C_{s,d}^{\prime S})}{C_{10\,s,d}^{SM}} \sqrt{1 - \frac{4m_\mu^2}{m_{B_{s,d}}^2}},$$

$$P_{s,d} \equiv \frac{m_{B_{s,d}}}{2m_\mu} \frac{(C_{s,d}^P - C_{s,d}^{\prime P})}{C_{10\,s,d}^{SM}} + \frac{(C_{s,d}^{10} - C_{10\,s,d}^{\prime})}{C_{10\,s,d}^{SM}}.$$

The $C_i$ are the Wilson coefficients corresponding to the Lagrangian

$$\mathcal{L}_s = \frac{4G_F}{\sqrt{2}} K_{tb} K_{ts}^* \frac{e^2}{16\pi^2} \sum_i (C_i O_i + C_i' O_i') + \mathrm{h.c.}\,.$$

[††]W. Altmannshofer and D.M. Straub, JHEP **1208**, 121 (2012) [arXiv:1206.0273 [hep-ph]].

The relevant operators for the $B_s$ decay are

$$O_s^{(\prime)S} = \frac{m_b}{m_{B_s}}(\bar{s}P_{R(L)}b)(\bar{\ell}\ell),$$
$$O_s^{(\prime)P} = \frac{m_b}{m_{B_s}}(\bar{s}P_{R(L)}b)(\bar{\ell}\gamma^5\ell),$$
$$O^{(\prime)}_{10\,s} = (\bar{s}\gamma_\mu P_{L(R)}b)(\bar{\ell}\gamma^\mu\gamma^5\ell),$$

The heavy Higgs $s$-channel tree-level diagrams contributing to $B_s$ decay yield

$$C_s^P = -Z_s\frac{m_{B_s}}{m_b}\frac{\rho_{32}^{D*}}{\sqrt{2}}\frac{m_\mu}{v}\tan\beta\frac{1}{m_A^2},$$
$$C_s^{\prime P} = Z_s\frac{m_{B_s}}{m_b}\frac{\rho_{23}^{D}}{\sqrt{2}}\frac{m_\mu}{v}\tan\beta\frac{1}{m_A^2} \ll C_s^P,$$
$$C_s^S = -Z_s\frac{m_{B_s}}{m_b}\sin(\beta-\alpha)\frac{\rho_{32}^{D*}}{\sqrt{2}}\frac{m_\mu}{v}\frac{\cos\alpha}{\cos\beta}\frac{1}{m_H^2},$$
$$C_s^{\prime S} = -Z_s\frac{m_{B_s}}{m_b}\sin(\beta-\alpha)\frac{\rho_{23}^{D}}{\sqrt{2}}\frac{m_\mu}{v}\frac{\cos\alpha}{\cos\beta}\frac{1}{m_H^2} \ll C_s^S,$$

where $Z_s \equiv \frac{16\pi^2\sqrt{2}}{4G_F K_{tb}K_{ts}^* e^2}$. Similar expressions are obtained for $B_d$ decay.

For the SM prediction, we take $C^{SM}_{10\,s,d} = -4.1$ and[‡‡]

$$\mathrm{BR}(B_s \to \mu^+\mu^-)_{SM} = (3.65 \pm 0.23) \times 10^{-9}\,,$$

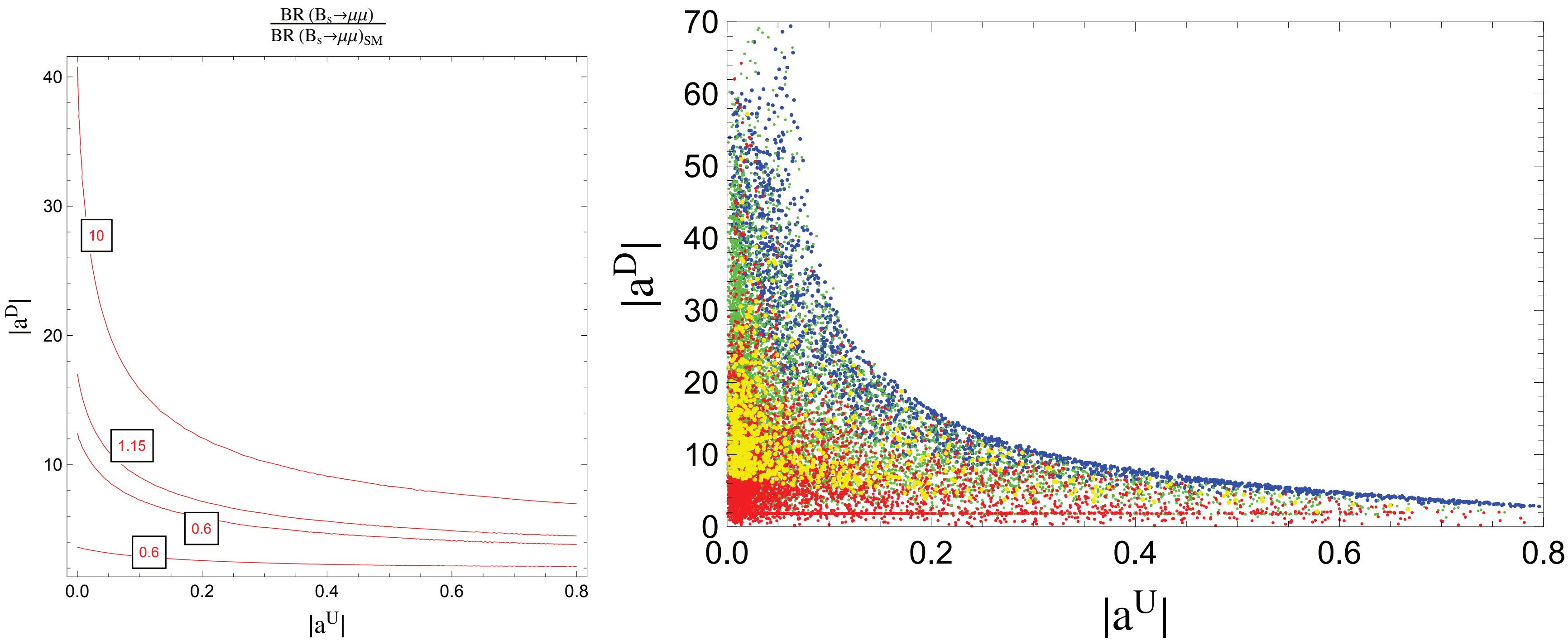
These values are in good agreement with the combination of the LHCb and the CMS measurements at Run I for the $B_s$ decay, which yields

$$\mathrm{BR}(B_s \to \mu^+\mu^-)_{\mathrm{exp}} = (2.8^{+0.7}_{-0.6}) \times 10^{-9}\,.$$

In what follows, we shall make use of the $2\sigma$ bound,

$$0.6 < \frac{\mathrm{BR}(B_s \to \mu^+\mu^-)_{\mathrm{exp}}}{\mathrm{BR}(B_s \to \mu^+\mu^-)_{\mathrm{SM}}} < 1.15\,.$$

[‡‡]C. Bobeth, M. Gorbahn, T. Hermann, M. Misiak, E. Stamou and M. Steinhauser, Phys. Rev. Lett. **112**, 101801 (2014) [arXiv:1311.0903 [hep-ph]].

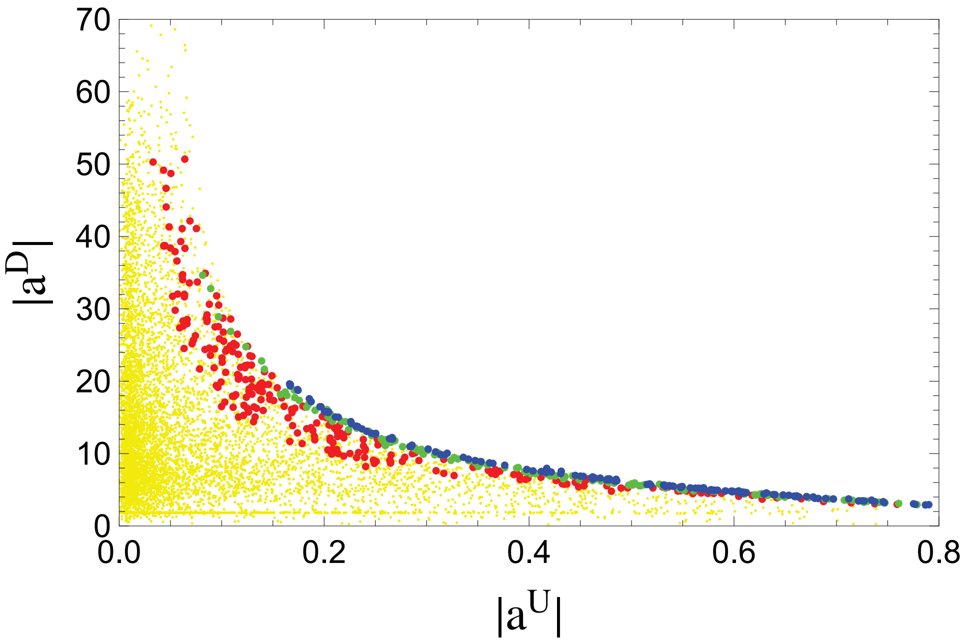



Left: The leading log approximation for BR$(B_s \to \mu^+\mu^-)$/BR$(B_s \to \mu^+\mu^-)_{\rm SM}$. Right: the same but scanning the parameter space. Yellow, red, green and blue colors correspond to a ratio of branching ratios of $[0.4, 0.6]$, $[0.6, 1.15]$, $[1.15, 10]$, and $> 10$. We have fixed $\beta - \alpha = \pi/2 - 0.2$ and $\Lambda_H = 400$ GeV.

3. $H \to b\bar{s},\ \bar{b}s$

$$\Gamma(H \to \bar{f}_i f_j) = \frac{3G_F v^2}{16\sqrt{2}\pi} m_H s^2_{\alpha-\beta}(|\rho^F_{ij}|^2 + |\rho^F_{ji}|^2)$$

$$\times \left[1 - \left(\frac{m_{f_i} - m_{f_j}}{m_H}\right)^2\right] \left[\left(1 - \frac{m^2_{f_i} + m^2_{f_j}}{m^2_H}\right)^2 - \frac{4m^2_{f_i} m^2_{f_j}}{m^4_H}\right]^{1/2} \quad (i \neq j)\,.$$




Yellow, red, green and blue colors correspond to $\mathrm{BR}(H \to bs)$ of $< 5 \times 10^{-4}$, $[5 \times 10^{-4}, 0.01]$, $[0.01, 0.1]$, and $> 0.1$ based on a full numerical scan. We have fixed $\beta - \alpha = \pi/2 - 0.2$ and $\Lambda_H = 400$ GeV.

# Conclusions

- In the search for new Higgs bosons, one should try to make the minimal set of assumptions that are consistent with the observed Higgs data.
- Current electroweak and Higgs data suggest a SM-like Higgs boson and highly suppressed FCNCs mediated by tree-level neutral Higgs exchange.
- Although special forms of the Higgs-fermion Yukawa couplings can naturally suppress FCNCs, one can imagine a more general set of assumptions that yield sufficiently suppressed Higgs-mediated FCNCs.
- In this talk, a framework was considered in which there is flavor alignment at a very high energy scale, which induces small Higgs-mediated FCNCs at the electroweak scale that can be consistent with current data.
- Some phenomenological consequences were examined, with an emphasis on processes that can distinguish among different models for the flavor structure of Higgs-fermion Yukawa interactions.